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D. Ebert and R. N. Faustov

Institut für Physik, Humboldt–Universität zu Berlin,

Invalidenstr.110, D-10115 Berlin

V. O. Galkin

Russian Academy of Sciences, Scientific Council for Cybernetics,

Vavilov Street 40, Moscow 117333, Russia

A. P. Martynenko and V. A. Saleev

Samara State University, Samara 443011, Russia

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Abstract

In the framework of the relativistic quasipotential quark model the mass spectrum of baryons with two heavy quarks is calculated. The quasipotentials for interactions of two quarks and of a quark with a scalar and axial vector diquark are evaluated. The bound state masses of baryons with $J^P = \frac{1}{2}^+, \frac{3}{2}^+$ are computed.

1 Introduction

The investigation of properties of hadrons, containing heavy quarks, is of a great interest for understanding the dynamics of quark and gluon interaction. To calculate the meson and baryon mass spectra QCD sum rules [1], potential models [2, 3, 4, 5], lattice QCD [6], heavy quark effective theory (HQET) [7, 8], the method of vacuum correlators [9] and some other approaches are widely used. Presently at the LHC, B-factories and the Tevatron with high luminosity, several experiments have been proposed, in which a detailed study of baryons containing two heavy quarks can be performed. The possible quark composition of such baryons looks as follows: (ccq, cbq, bbq), where c and b are

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†On leave of absence from Russian Academy of Sciences, Scientific Council for Cybernetics, Vavilov Street 40, Moscow 117333, Russia.

heavy quarks (Q) and q denotes a light u, d and s quark. In this connection theoretical predictions for heavy baryon masses acquire important significance. The properties of heavy hadrons (containing b and/or c quarks) are essentially different from those of light hadrons composed from u, d and s quarks. This difference originates from the fact that the heavy quark mass strongly exceeds the scale of QCD interaction, $\Lambda_{QCD} \approx 300-400$ MeV, namely $m_{c,b} \gg \Lambda_{QCD}$. Therefore, while investigating processes with heavy quarks one should expect that the limit $m_Q \to \infty$ within QCD yields good quantitative results, which may be further improved in HQET by taking into account $1/m_Q$ corrections. Baryons, containing two heavy quarks (QQq), may be considered as a localized source (QQ) of the colour field, in which the light quark moves. Interaction forces between heavy quarks then lead to the formation of a two-particle bound state of a QQ-diquark, the scale of which is determined by the quantity $1/m_Q$ which is small compared to the QCD scale $1/\Lambda_{QCD}$. Thereby it seems justified to treat the heavy diquark as a pointlike object with definite colour, spin and mass and to ignore the configuration with a qQ-diquark which size is of order $1/m_q \sim 1/\Lambda_{QCD}$. The baryon formation then occurs as a result of a diquark interaction with the light quark. Here in the framework of the quark-diquark approximation we calculate the mass spectrum of heavy baryons on the basis of a local Schrödinger-like quasipotential equation [10]:

$$\left(\frac{b^2(M)}{2\mu_R} - \frac{\vec{p}^2}{2\mu_R}\right)\psi_M(\vec{p}) = \int \frac{d^3q}{(2\pi)^3} V(\vec{p},\vec{q},M)\psi_M(\vec{q}), \tag{1}$$

where the relativistic reduced mass is

$$\mu_R = \frac{E_1E_2}{E_1+E_2} = \frac{M^4-(m_1^2-m_2^2)^2}{4M^3}, \tag{2}$$

$$E_1 = \frac{M^2-m_2^2+m_1^2}{2M}, \quad E_2 = \frac{M^2-m_1^2+m_2^2}{2M}, \quad E_1+E_2 = M,$$

and the center of mass system relative momentum squared on the mass shell reads

$$b^2(M) = \frac{[M^2-(m_1+m_2)^2][M^2-(m_1-m_2)^2]}{4M^2}, \tag{3}$$

with $m_{1,2}$ the masses of the constituent particles. In refs. [11, 12] the operator of quark-antiquark interaction has been constructed and the mass spectra and decay rates of mesons have been investigated. The relativistic quasipotential quark model gives in the meson sector results which nicely agree with experimental data. Relativistic effects play an important role in describing properties of quark bound states and may be consistently taken into account within the quasipotential method [13].

2 Diquarks in the relativistic quasipotential model

The concept of diquarks was introduced at the very beginning of the development of the quark model of hadrons in order to describe baryon properties [5]. Diquarks represent

by themselves two-particle clusters, which are produced inside the three-quark systems as a result of quark interactions in which the spin-dependent forces play an important role [14, 15]. The attractive forces between two quarks act if the quarks are in the antisymmetric colour state. The interaction potential at large distances in the quark-quark system is similar to that in the quark-antiquark system since the baryon size is practically the same as the meson size. In the ground state the diquarks are two-particle bound states of quarks in an antisymmetric colour state with zero angular momentum and definite flavour and spin. In the case of indentical quarks the diquark has spin $S = 1$, whereas for quarks of different flavours the state with spin $S = 0$ is also possible.

Now we construct the interaction quasipotential for two heavy quarks Q, which enters in eq. (1). The interaction operator in the QQ system contains a perturbative part, which is defined at small distances in the quasipotential approach by the one-gluon exchange amplitude and can be evaluated within QCD. Moreover there is a nonperturbative term in the operator $V(\vec{p},\vec{q},M)$, which cannot be obtained consistently within QCD and leads to the quark confinement at large distances.

The quasipotential $V(\vec{p},\vec{q},M)$ of eq. (1) is constructed with the help of the off-mass-shell scattering amplitude, projected onto the positive energy states. Then the perturbative part of the quasipotential can be represented in the form:

$$V^{pert.}_{QQ}(\vec{p},\vec{q}) = \bar{u}_1(\vec{p})\bar{u}_2(-\vec{p})\frac{\frac{2}{3}\alpha_s}{4\sqrt{\epsilon_1(\vec{p})\epsilon_1(\vec{q})\epsilon_2(\vec{p})\epsilon_2(\vec{q})}}D_{\mu\nu}(k)\gamma_1^{\mu}\gamma_2^{\nu}u_1(\vec{q})u_2(-\vec{q}), \tag{4}$$

where $D_{\mu\nu}(k)$ is the gluon propagator conveniently taken in the Coulomb gauge

$$D^{00}(k) = -\frac{1}{\vec{k}^2}, \quad D^{ij}(k) = -\frac{1}{k^2}\left(\delta^{ij} - \frac{k^ik^j}{\vec{k}^2}\right), \quad D^{0i} = D^{i0} = 0, \tag{5}$$

$\vec{k} = \vec{p} - \vec{q}$, and $u(\vec{p})$ is the Dirac spinor,

$$u^{\lambda}(\vec{p}) = \sqrt{\varepsilon(\vec{p}) + m}\begin{pmatrix} 1 \\ \frac{\vec{\sigma}\vec{p}}{\varepsilon(\vec{p})+m}\end{pmatrix}\chi^{\lambda}. \tag{6}$$

In this case the colour factor, which includes the colour part of the baryon wave function $\varepsilon_{ijk}/\sqrt{6}$ $(i,j,k = 1,2,3)$ has the value:

$$\frac{1}{\sqrt{6}}\varepsilon_{ijk}\delta_{nk}\frac{1}{\sqrt{6}}\varepsilon_{lmn}T^a_{li}T^a_{mj} = -\frac{2}{3}, \tag{7}$$

where T^a are the generators of the colour $SU_c(3)$ group. It means, that in the two-quark system the exchange of a nonabelian gauge boson results in an attractive Coulomb potential at small distances which differs only by a factor 1/2 from the similar potential in the quark-antiquark system. Substituting the spinor (6) into eq.(4) and carrying out some manipulations with spinor convolutions, we obtain, in the same way as in [12], the one-gluon exchange quasipotential including relativistic corrections. Since we consider the

ground state of diquarks, the angular momentum of the QQ system is zero. As it has been remarked in ref. [16], the quark spin interaction can result in the existence of short-range correlations in the quark-quark subsystem. Therefore, neglecting corrections $O(p^2/m_Q^2)$, we can represent the perturbative part of the quasipotential in the configuration space in the form:

$$V_{QQ}^{pert.}(r) = -\frac{2}{3}\frac{\alpha_s}{r} + \frac{4\pi\alpha_s}{9m_1m_2}(\vec{\sigma}_1\vec{\sigma}_2)\delta(\vec{r}), \tag{8}$$

where $\alpha_s(r)$ is the running coupling constant and $\sigma_{1,2}$ are the quark spin matrices. The colour hyperfine interaction is attractive if the total spin of two quarks $S = 0$, and repulsive in the case $S = 1$.

In the approximation of one-gluon exchange the quark-quark potential in a baryon is equal to one half of the quark-antiquark potential in a meson. Calculations based on the baryon Wilson loop area law in QCD [17] show that this relation holds even true for nonperturbative interactions. Thus, while constructing the long-range nonperturbative part of the quark interaction operator we used the known "1/2 rule", i. e. $V_{qq} = \frac{1}{2}V_{q\bar{q}}$ [5] and chose this part of the quasipotential in the standard linearly rising in r form. As a result in this approximation the complete interaction operator in the quark-quark system is chosen in the form:

$$V_{QQ}(r) = -\frac{2}{3}\frac{\alpha_s}{r} + \frac{1}{2}(Ar + B) + \frac{4\pi\alpha_s}{9m_1m_2}(\vec{\sigma}_1\vec{\sigma}_2)\delta(\vec{r}), \tag{9}$$

where the parameters $A = 0.18$ GeV^2 and $B = -0.3$ GeV had been fixed previously [11, 12] in calculating the meson mass spectrum. The numerical solution of the quasipotential equation (1) with interaction operator (9) yields the mass spectrum of the s-wave axial vector and scalar diquarks. The results are presented [1] in Table 1. Here the usual values of quark masses (the same as for mesons) were used [11, 12]:

$m_u = m_d = 0.33$ GeV, $m_s = 0.5$ GeV, $m_c = 1.55$ GeV, $m_b = 4.88$ GeV.

Table 1: Masses of s-wave scalar and axial vector diquarks

Diquark	qq	qs	qc	qb	ss	sc	sb	cc	cb	bb
M(GeV) (scalar diquark)	0.85	1.04	2.08	5.39	-	2.27	5.57	-	6.52	-
M(GeV) (axial vector diquark)	1.02	1.15	2.12	5.41	1.32	2.29	5.58	3.26	6.52	9.79

[1] The masses of light diquarks are presented for completeness where $q = u, d$.

3 Quasipotential for the interaction of a quark with a scalar diquark

Now to calculate the baryon mass spectrum in the quark-diquark approximation it is necessary to construct the quark-diquark interaction operator. In the case of a scalar diquark the perturbative part of the quasipotential is described by the Feynman diagram in Fig. 1.

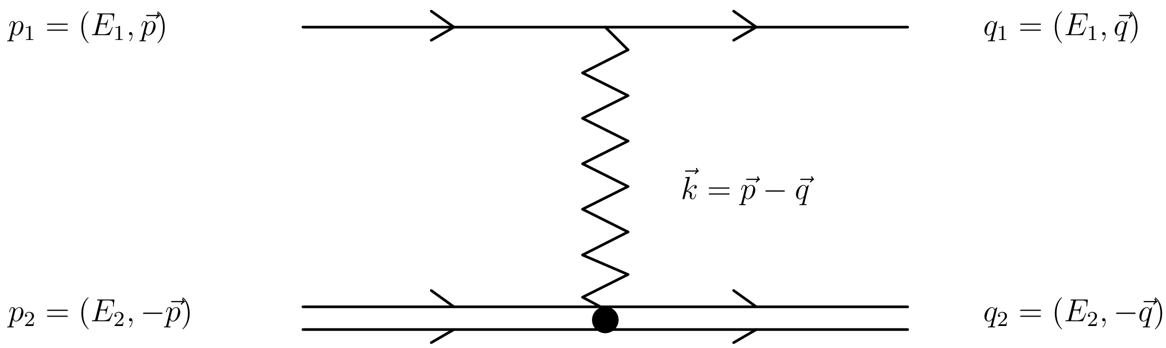

Figure 1: Feynman diagram of the one-gluon interaction in the quark-diquark system.

The corresponding expression for the interaction operator has the form:

$$V^{pert.}_{q+SD}(\vec{p},\vec{q}) = \frac{\frac{4}{3}g_s^2}{4\sqrt{\epsilon_1(\vec{p})\epsilon_1(\vec{q})\epsilon_2(\vec{p})\epsilon_2(\vec{q})}}(p_2+q_2)^\mu D_{\mu\nu}(\vec{k})\bar{u}(\vec{q})\gamma^\nu u(\vec{p}), \tag{10}$$

where the interaction vertex of a scalar diquark with a gluon is given by the factor $ig_sT^a_{ij}(p_2+q_2)^\mu$. Using the standard representation of the Dirac spinors (6), we get for $V^{pert.}_{q+SD}$ up to terms $O(\vec{p}^2/m^2)$ the following expression:

$$V^{pert.}_{q+SD}(\vec{p},\vec{q}) = -\frac{\frac{4}{3}g_s^2 4\pi}{\vec{k}^2}\Bigg\{\left[1+\frac{b^2}{E_1E_2}\right] + \frac{\vec{p}^2-b^2}{4m_1^2} + \frac{\vec{p}^2-b^2}{4m_2^2} + \frac{\vec{p}^2-b^2}{m_1m_2} - \tag{11}$$

$$-(\vec{p}\vec{k})\left(\frac{1}{4m_1^2}+\frac{1}{4m_2^2}+\frac{1}{m_1m_2}\right) + \frac{\vec{k}^2}{4m_1m_2} + i\vec{S}_1[\vec{p}\times\vec{q}]\frac{1}{2m_1}\left(\frac{1}{m_1}+\frac{2}{m_2}\right)\Bigg\}.$$

Then we add to this expression in the configuration space the confining potential

$$V^{nonpert.}_{q+SD}(r) = Ar + B, \tag{12}$$

which is the same as in the quark-antiquark system. Thus the complete quasipotential, which is used to calculate the mass spectrum of baryons with $J^P = \frac{1}{2}^+$ can be represented as follows:

$$V^{tot.}_{q+SD}(r) = Ar + B - \frac{4\alpha_s}{3r}\left(1+\frac{b^2}{E_1E_2}\right) - \frac{4\pi\alpha_s}{3m_1m_2}\delta(\vec{r}) + \frac{4\alpha_s}{3m_1m_2}\left(\frac{m_2}{2m_1}+1\right)\frac{\vec{L}\vec{S}_1}{r^3} - \tag{13}$$

$$-\frac{2\mu\alpha_s^2}{3r^2}\left(\frac{1}{m_1^2}+\frac{1}{m_2^2}+\frac{4}{m_1m_2}\right)-\frac{4\alpha_s}{3r^3}(\vec{r}\nabla)\left(\frac{1}{4m_1^2}+\frac{1}{4m_2^2}+\frac{1}{m_1m_2}\right),$$

where m_2 is the scalar diquark mass, m_1 is the light quark mass and μ is their reduced mass. The numerical results for the mass spectrum of the ground states of baryons with $J^P=\frac{1}{2}^+$ on the basis of eqs. (1) and (13) are given in Table 2.

4 Quasipotential for the interaction of a quark with an axial vector diquark

As in the previous section, the perturbative part of the interaction operator for spin 1/2 and spin 1 particles is given by the diagram in Fig. 1. It includes the matrix element of the chromomagnetic current operator between states of the axial vector diquark. This matrix element is defined by three chromomagnetic form factors. Neglecting the chromo-quadrupole moment of the axial vector particle, we can represent the four vector of the chromomagnetic current in the form [18]:

$$j^\mu(\vec{p},\vec{q})=igT^a_{ij}\xi^*\Big\{G_1(k^2)(p_2+q_2)^\mu+\frac{G_2(k^2)}{m_2}[W^\mu(\vec{p}_2)(\vec{S}_2\vec{\Delta}_2)-(\vec{S}_2\vec{\Delta}_2)W^\mu(\vec{p}_2)]\Big\}\xi, \quad (14)$$

$$\vec{\Delta}_2=-\vec{q}+\frac{\vec{p}}{m_2}\left(E_2-\frac{\vec{q}\vec{p}}{E_2+m_2}\right),$$

where $\vec{S}_2$ is the spin operator of a vector particle, W^μ is the four vector of the relativistic spin (the Pauli-Lubansky vector) with the components:

$$W^0(\vec{p}_2)=-\vec{S}_2\vec{p}, \qquad \vec{W}(\vec{p}_2)=m_2\vec{S}_2+\frac{\vec{p}(\vec{S}_2\vec{p})}{E_2+m_2}, \quad (15)$$

and ξ is the axial vector particle wave function (polarization vector) in its rest frame. Then the one-gluon exchange quasipotential takes the form:

$$V^{pert.}_{q+AD}(\vec{p},\vec{q})=\frac{\frac{4}{3}g_s^2}{4\sqrt{\epsilon_1(\vec{p})\epsilon_1(\vec{q})\epsilon_2(\vec{p})\epsilon_2(\vec{q})}}D_{\mu\nu}(\vec{k})\bar{u}(\vec{q})\gamma^\nu u(\vec{p}) \quad (16)$$

$$\times\Big[G_1(k^2)(p_2+q_2)^\mu+G_2(k^2)[\frac{1}{m_2}W^\mu(\vec{p}_2)(\vec{S}_2\vec{\Delta}_2)-\frac{1}{m_2}(\vec{S}_2\vec{\Delta}_2)W^\mu(\vec{p}_2)]\Big].$$

Now we expand eq. (16) with an accuracy up to terms $O(\vec{p}^2/m^2)$. The zeroth and spacial vector components of the current (14) contain the following terms:

$$\frac{1}{m_2}W^0(\vec{p}_2)(\vec{S}_2\vec{\Delta}_2)-\frac{1}{m_2}(\vec{S}_2\vec{\Delta}_2)W^0(\vec{p}_2)\cong\frac{i}{m_2}\vec{S}_2[\vec{p}\times\vec{q}], \quad (17)$$

$$\frac{1}{m_2}\vec{W}(\vec{p}_2)(\vec{S}_2\vec{\Delta}_2)-\frac{1}{m_2}(\vec{S}_2\vec{\Delta}_2)\vec{W}(\vec{p}_2)\cong i[(\vec{p}-\vec{q})\times\vec{S}_2], \quad (18)$$

which produce fine and hyperfine splittings of energy levels. The chromomagnetic form factor of the axial vector particle $G_2(k^2)$ at $k^2 = 0$ defines its total chromomagnetic moment: $G_2(0) = \mu_V = 1 + \zeta$. Setting here the anomalous chromomagnetic moment of the axial vector diquark $\zeta = 1$ and $G_1(0) = 1$, we obtain the following expression for $V^{pert.}_{q+AD}$:

$$V^{pert.}_{q+AD}(\vec{p},\vec{q}) = -\frac{4}{3}\frac{g_s^2}{\vec{k}^2}\frac{4\pi}{}\left\{\left(1+\frac{b^2}{E_1E_2}\right)+\frac{\vec{p}^2-b^2}{4m_1^2}+\frac{\vec{p}^2-b^2}{4m_2^2}+\frac{\vec{p}^2-b^2}{m_1m_2}-\right. \tag{19}$$
$$-(\vec{p}\vec{k})(\frac{1}{4m_1^2}+\frac{1}{4m_2^2}+\frac{1}{m_1m_2})+\frac{\vec{k}^2}{4m_1m_2}+i\vec{S}_1[\vec{p}\times\vec{q}]\frac{1}{2m_1}\left(\frac{1}{m_1}+\frac{2}{m_2}\right)+$$
$$\left.+i\vec{S}_2[\vec{p}\times\vec{q}]\frac{1}{2m_2}\left(\frac{1}{m_2}+\frac{2}{m_1}\right)-\frac{1}{m_1m_2}[(\vec{S}_1\vec{S}_2)\vec{k}^2-(\vec{S}_1\vec{k})(\vec{S}_2\vec{k})]\right\}.$$

The nonperturbative part of the interaction potential for a quark and an axial vector diquark is taken in the form (12). Transforming (19) into configuration space, we get the complete expression of the interaction operator for the system $q + AD$:

$$V^{tot.}_{q+AD}(r) = Ar + B - \frac{4\alpha_s}{3r}\left(1+\frac{b^2}{E_1E_2}\right)-\frac{4\pi\alpha_s}{3m_1m_2}\delta(\vec{r})+\frac{4\alpha_s}{3m_1m_2}\left(\frac{m_2}{2m_1}+1\right)\frac{\vec{L}\vec{S}_1}{r^3}- \tag{20}$$
$$-\frac{2\mu\alpha_s^2}{3r^2}\left(\frac{1}{m_1^2}+\frac{1}{m_2^2}+\frac{4}{m_1m_2}\right)-\frac{4\alpha_s}{3r^3}(\vec{r}\nabla)\left(\frac{1}{4m_1^2}+\frac{1}{4m_2^2}+\frac{1}{m_1m_2}\right)+$$
$$+\frac{4\alpha_s}{3m_1m_2}\left(\frac{m_1}{2m_2}+1\right)\frac{\vec{L}\vec{S}_2}{r^3}+\frac{32\pi\alpha_s}{9m_1m_2}(\vec{S}_1\vec{S}_2)\delta(\vec{r})+\frac{4\alpha_s}{3m_1m_2r^3}[(\vec{S}_1\vec{S}_2)-\frac{3}{r^2}(\vec{S}_1\vec{r})(\vec{S}_2\vec{r})].$$

The potential (20) has been used for calculating the mass spectrum of baryons with $J^P = \frac{1}{2}^+, \frac{3}{2}^+$ on the basis of the local quasipotential equation (1). Corresponding numerical results are presented in Table 2.

5 Conclusions

In this paper we have carried out the calculation of the mass spectrum of baryons with two heavy quarks in the quark-diquark approximation on the basis of a local quasipotential equation. As it follows from Table 2, our results are in good agreement with calculations made within other approaches. One can see from Table 2 that our predictions are especially close to those obtained in [19]. The approach used in [19, 20] is based on some semiempirical regularities of the baryon mass spectrum following from experimental data and the constituent quark model. The authors of [19, 20] found a successful parametrization of the interpolating curves which describe the known meson and baryon masses, and give predictions for some yet unobserved mesons and baryons. Our results for the B_c meson mass spectrum [12] are also rather close to those of [20]. Our only remark is that one should take into account the mixing of $^3P_1$ and $^1P_1$ B_c states. The quoted authors use in

Table 2: Masses of baryons (in GeV) containing two heavy quarks. $[QQ]$ denotes the diquark subsystem with the antisymmetric spin wave function, $\{QQ\}$ denotes the diquark subsystem with the symmetric spin wave function ($q = u, d$).

Notation	Quark content	J^P	M_B (our results)	M_B [19]	M_B [8]
Ξ_{cc}	$\{cc\}q$	$\frac{1}{2}^+$	3.66	3.66	3.61
Ξ^*_{cc}	$\{cc\}q$	$\frac{3}{2}^+$	3.81	3.74	3.68
Ω_{cc}	$\{cc\}s$	$\frac{1}{2}^+$	3.76	3.74	3.71
Ω^*_{cc}	$\{cc\}s$	$\frac{3}{2}^+$	3.89	3.82	3.76
Ξ_{bb}	$\{bb\}q$	$\frac{1}{2}^+$	10.23	10.34	-
Ξ^*_{bb}	$\{bb\}q$	$\frac{3}{2}^+$	10.28	10.37	-
Ω_{bb}	$\{bb\}s$	$\frac{1}{2}^+$	10.32	10.37	-
Ω^*_{bb}	$\{bb\}s$	$\frac{3}{2}^+$	10.36	10.40	-
Ξ_{cb}	$\{cb\}q$	$\frac{1}{2}^+$	6.95	7.04	-
Ξ'_{cb}	$[cb]q$	$\frac{1}{2}^+$	7.00	6.99	-
Ξ^*_{cb}	$\{cb\}q$	$\frac{3}{2}^+$	7.02	7.06	-
Ω_{cb}	$\{cb\}s$	$\frac{1}{2}^+$	7.05	7.09	-
Ω'_{cb}	$[cb]s$	$\frac{1}{2}^+$	7.09	7.06	-
Ω^*_{cb}	$\{cb\}s$	$\frac{3}{2}^+$	7.11	7.12	-

particular the Feynman-Hellmann theorem in order to obtain the monotonic decrease of the hadron mass spectrum with respect to quark masses. Here they use an oversimplified assumption about the interquark potential, namely the complete neglection of the spin-independent part of the Breit-Fermi potential as being unimportant. This assumption leads finally to some inequalities between quark masses and meson masses. These inequalities are violated in many constituent quark models (including ours) in which widely adopted quark masses are used. From our point of view there is nothing serious in these contradictions and the recovery of the appropriate spin-independent (flavour-dependent) part of the quark potential [12] will eliminate them. Thus we may conclude that our model correctly describes some important features of the baryon mass spectrum.

There exists yet another type of baryons with one heavy and two light quarks for which the quark-diquark approximation gives reasonable results. Thus, in Ref. [21] Qqq baryons were treated by considering the interaction of the diquark system with the third quark on the basis of additional quark exchange forces [22, 23, 24]. The resulting effective lagrangian, which incorporates heavy quark and chiral symmetry, describes interactions of heavy baryons with Goldstone bosons in the low energy region. As an interesting application, the Isgur-Wise form factors have been calculated.

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